

Branching Fractions and CP Asymmetries of $B \rightarrow K_0^*(1430)\rho$ and $B \rightarrow K_0^*(1430)\phi$ Decays in the Family Nonuniversal Z' Model

Ying Li* and En-Lei Wang

Department of Physics, Yantai University, Yantai 264-005, China

(Dated: March 8, 2013)

In this work, within the QCD factorization approach, we investigate the branching fractions and CP asymmetries of decays $B \rightarrow K_0^*(1430)\rho$ and $B \rightarrow K_0^*(1430)\phi$ under two different scenarios both in the standard model and the family nonuniversal Z' model. We find that the annihilation terms play crucial roles in these decays and lead to the main uncertainties. For decays $B^- \rightarrow \bar{K}_0^{*-}(1430)\rho^0(\omega)$, the new Z' boson could change branching fractions remarkably. However, for other decays, its contribution might be clouded by large uncertainties from annihilations. Unfortunately, neither the standard model nor Z' model can reproduce all experimental data under one certain scenario. We also noted that the CP asymmetries of $B^- \rightarrow \bar{K}_0^{*-}(1430)\rho^0(\omega)$ could be used to identify the $K_0^*(1430)$ meson and search for the new physics contribution.

I. INTRODUCTION

The study of B meson rare decays is a crucial tool in testing the fundamental interactions among elementary particles, exploring the origin of CP violation, and searching for possible new physics (NP) beyond the standard model (SM). Theoretically and experimentally, such kind of research has been conducted in great detail, especially in the weak interactions of B meson. In particular, the processes induced by flavor-changing neutral-current (FCNC) only occur at the loop level in SM, and are therefore a very sensitive probe of NP beyond SM. Already, FCNC processes have been explored mainly in the $B_q - \bar{B}_q$ mixing and the semi-leptonic weak decays, which permit a clean theoretical description. So far, the charmless hadronic B meson decays induced by FCNC have also been studied extensively, such as $B \rightarrow K\pi, K^{(*)}\phi$ and $K^{(*)}\eta^{(\prime)}$ decays. In the past few years, the new physics effect in these decays have also been studied widely, such as in supersymmetry

*Electronic address: liying@ytu.edu.cn

model, two-Higgs doublet model, Z' model, the forth generation model, extra dimension models, and so on (see review in [1] and references therein).

In order to search for effect of NP in the nonleptonic B decays, most theoretical studies are focused on $B \rightarrow PP, PV$ or VV in the past few years. But, the studies of decay modes involving a scalar meson are relatively few, because the underlying structure of the scalar mesons is not well established in theoretical side. To describe the component of the scalar mesons, there are usually two possible scenarios (S1 and S2) according to the QCD sum rule method [2]: (i) In S1, we treat scalars above 1 GeV as the first excited states, while the scalars under 1 GeV are regarded as the low lying states; (ii) In S2, the scalars above 1 GeV are viewed as the ground states, and light scalars are four-quark bound states or hybrid states. Under these two scenarios, many special decays have been examined within the QCD factorization (QCDF) approach [3, 4] or the perturbative QCD approach (pQCD) [5–14]. However, because of large uncertainties in SM, the NP effects in these decays are rarely studied.

Very recently, BaBar collaboration reported their first branching fraction measurements for the decays $B \rightarrow K_0^*(1430)\rho$ that are induced by FCNC [15]:

$$Br(B^0 \rightarrow K_0^{*0}(1430)\rho^0) = (27 \pm 4 \pm 2 \pm 3) \times 10^{-6}; \quad (1)$$

$$Br(B^0 \rightarrow K_0^{*+}(1430)\rho^-) = (28 \pm 10 \pm 5 \pm 3) \times 10^{-6}. \quad (2)$$

The above results are inconsistent with the pQCD predictions [12] in most cases. Moreover, these results are somewhat much lower than the QCDF predictions [4] but are consistent with QCDF within rather large uncertainties. For $B \rightarrow K_0^*(1430)\phi$, BaBar collaborator also updated their results [16, 17] in ref. [18]:

$$Br(B^0 \rightarrow K_0^{*0}(1430)\phi) = (4.3 \pm 0.6 \pm 0.4) \times 10^{-6}; \quad (3)$$

$$Br(B^\pm \rightarrow K_0^{*\pm}(1430)\phi) = (7.0 \pm 1.3 \pm 0.9) \times 10^{-6}. \quad (4)$$

Both QCDF and pQCD calculation of above modes have also been presented in Refs.[4, 8], and the predicted central values of $B^0 \rightarrow \phi K_0^{*0}(1430)$ deviate from the experimental data, though they can be also accommodated within very large theoretical errors. In the following, $K_0^*(1430)$ is denoted as K_0^* in some places for convenience.

The predictions of SM cannot agree the data convincingly, which gives us possible hints on physics beyond SM. It is our purpose of this work to show that a new physics effect of similar size can be obtained from some models with an extra spin-1 Z' bosons, which are known to naturally

exist in some well-motivated extensions of the SM [19]. Interesting phenomena arise when the Z' couplings to physical fermion eigenstates are nondiagonal, which could be realized in the E_6 models [20], string models [21] and some grand unified theories [22]. For example, in the super string model advocated by Chaudhuri *et.al.*[21], it is possible to have family nonuniversal Z' couplings, because of different constructions of the different families. It also should be note that in such a model, called the family nonuniversal Z' model, the nonuniversal couplings could lead to FCNCs at the tree level as well as introduce new weak phases [23], which could explain the CP asymmetries in the current high energy experiments. In fact, the effects of Z' models have been studied extensively in the low energy flavor physics phenomena, such as neutral mesons mixing, B meson decays, single top production and lepton decays [23–31].

In this current work, we shall adopt the QCD factorization approach [32] to evaluate the relevant hadronic matrix elements of B decays, since it is a systematic framework to calculate these matrix elements from QCD theory, and holds in the heavy quark limit $m_b \rightarrow \infty$ and the heavy quark symmetry. In such calculations, one requires the additional knowledge about form factors of B meson to the scalar or the vector transitions. This problem, being a part of the nonperturbative sector of QCD, lacks a precise solution. To the best of our knowledge, a number of different approaches had been used to calculate the form factors of $B \rightarrow S$ decays, such as QCD sum rule [33, 34], light-cone QCD sum rule [35, 36], perturbative QCD approach [37] and covariant light front quark model (cLFQM) [38]. Among them, the form factors of the cLFQM are first calculated in the spacelike region and their momentum dependence is fitted to a 3-parameter form. This parameterization is then analytically continued to the timelike region to determine the physical form factors at $q^2 \geq 0$. Moreover, for these form factors both the heavy quark limit and heavy quark symmetry are satisfied. For that reason, we will use the results of cLFQM [38] in the following calculations.

For comparison, $B \rightarrow K_0^* \rho$ and $K_0^* \phi$ decays in SM should be reinvestigated in Section.II. In Section.III, we will review the family nonuniversal Z' model briefly and show the effect of Z' to decay modes we are considering. In Section.IV, we will present our numerical results and discussions in great detail. At last, we will summarize this work in Section.V .

II. REVISITING $B \rightarrow \rho K_0^*(1430)$ AND $B \rightarrow \phi K_0^*(1430)$ DECAYS WITHIN THE QCDF FRAMEWORK

To proceed, we discuss the decay constants of the scalar meson. Unlike pseudoscalar meson, each scalar meson has two decay constants, the vector decay constant f_S and the scale-dependent scalar decay constant \bar{f}_S namely, which are defined as:

$$\langle S(p) | \bar{q}_2 \gamma_\mu q_1 | 0 \rangle = f_S p_\mu, \quad \langle S(p) | \bar{q}_2 q_1 | 0 \rangle = m_S \bar{f}_S, \quad (5)$$

and they are related by the equation of motion:

$$f_S = \frac{m_2(\mu) - m_1(\mu)}{m_S} \bar{f}_S, \quad (6)$$

where m_2 and m_1 are the running current quark masses. Therefore, the vector decay constant is much smaller than the scalar one. As for the vector meson, the two kinds of decay constants are also given by [39]

$$\langle V(p) | \bar{q}_2 \gamma_\mu q_1 | 0 \rangle = f_V m_V \epsilon_\mu^*, \quad \langle V(p, \epsilon^*) | \bar{q} \sigma_{\mu\nu} q' | 0 \rangle = f_V^\perp (p_\mu \epsilon_\nu^* - p_\nu \epsilon_\mu^*). \quad (7)$$

The twist-2 and twist-3 light-cone distribution amplitudes (LCDAs) of scalar mesons, $\phi_S(x)$, $\phi_S^s(x)$ and $\phi_S^\sigma(x)$ respect the normalization conditions:

$$\int_0^1 dx \phi_S(x) = \frac{f_S}{2\sqrt{6}}, \quad \int_0^1 dx \phi_S^s(x) = \int_0^1 dx \phi_S^\sigma(x) = \frac{\bar{f}_S}{2\sqrt{6}}, \quad (8)$$

and $\phi_S^T(x) = \frac{1}{6} \frac{d}{dx} \phi_S^\sigma(x)$. The twist-2 LCDA can be expanded in the Gegenbauer polynomials:

$$\phi_S(x, \mu) = \frac{1}{\sqrt{6}} \bar{f}_S(\mu) 6x(1-x) \sum_{m=1}^{\infty} B_m(\mu) C_m^{3/2}(2x-1). \quad (9)$$

The decay constants and the Gegenbauer moments of the twist-2 wave function in two different scenarios have been studied explicitly in Refs. [3] using the QCD sum rule approach. As for the explicit form of the Gegenbauer moments for the twist-3 wave functions, there exist some uncertainties theoretically [40], thus we choice the asymptotic form for simplicity:

$$\phi_S^s = \frac{1}{\sqrt{6}} \bar{f}_f, \quad \phi_S^T = \frac{1}{\sqrt{6}} \bar{f}_S (1-2x). \quad (10)$$

For the vector mesons, the normalization for the twist-2 function Φ_V and the twist-3 function Φ_ν is given by

$$\int_0^1 dx \Phi_V(x) = f_V, \quad \int_0^1 dx \Phi_\nu(x) = 0, \quad (11)$$

where the definitions for $\Phi_V(x)$ can be found in [32]. The general expressions of these LCDAs read

$$\Phi_V(x, \mu) = 6x(1-x)f_V \left[1 + \sum_{n=1}^{\infty} \alpha_n^V(\mu) C_n^{3/2}(2x-1) \right], \quad (12)$$

and

$$\Phi_V(x, \mu) = 3f_V^\perp \left[2x-1 + \sum_{n=1}^{\infty} \alpha_{n,\perp}^V(\mu) P_{n+1}(2x-1) \right], \quad (13)$$

where $P_n(x)$ are the Legendre polynomials.

In the calculation, the most important nonperturbative parameters are form factors of $B \rightarrow S, V$ transitions, which are defined by [41]:

$$\begin{aligned} \langle V(p') | V_\mu | B(p) \rangle &= -\frac{1}{m_B + m_V} \varepsilon_{\mu\nu\alpha\beta} \varepsilon^{*\nu} P^\alpha q^\beta V^{BV}(q^2), \\ \langle V(p') | A_\mu | B(p) \rangle &= i \left\{ (m_B + m_V) \varepsilon_\mu^* A_1^{BV}(q^2) - \frac{\varepsilon^* \cdot P}{m_B + m_V} P_\mu A_2^{BV}(q^2) \right. \\ &\quad \left. - 2m_V \frac{\varepsilon^* \cdot P}{q^2} q_\mu [A_3^{BV}(q^2) - A_0^{BV}(q^2)] \right\}, \\ \langle S(p') | A_\mu | B(p) \rangle &= -i \left[\left(P_\mu - \frac{m_B^2 - m_S^2}{q^2} q_\mu \right) F_1^{BS}(q^2) + \frac{m_B^2 - m_S^2}{q^2} q_\mu F_0^{BS}(q^2) \right], \end{aligned} \quad (14)$$

with $P_\mu = (p + p')_\mu$, $q_\mu = (p - p')_\mu$.

To calculate the amplitudes, we start from the effective Hamiltonian responsible for $b \rightarrow s$ transitions, which is given by [42]

$$\begin{aligned} \mathcal{H}_{\text{eff}} &= \frac{G_F}{\sqrt{2}} \left[V_{ub} V_{us}^* (C_1 O_1^u + C_2 O_2^u) + V_{cb} V_{cs}^* (C_1 O_1^c + C_2 O_2^c) - V_{tb} V_{ts}^* \left(\sum_{i=3}^{10} C_i O_i \right. \right. \\ &\quad \left. \left. + C_{7\gamma} O_{7\gamma} + C_{8g} O_{8g} \right) \right] + \text{h.c.} \end{aligned} \quad (15)$$

In the above equation, $V_{qb} V_{qs}^*$ ($q = u, c, t$) represent for products of the Cabibbo-Kobayashi-Maskawa (CKM) matrix elements, C_i are the responding Wilson coefficients, and O_i are the relevant four-quark operators whose explicit forms could be found, for example, in Refs. [42].

We now turn to study the short-distance contributions within the QCDF approach, where the contribution of the nonperturbative sector is dominated by the form factors and the nonfactorizable impact in the hadronic matrix elements is controlled by hard gluon exchange. The hadronic matrix

elements of the decay can be written as

$$\begin{aligned}\langle M_1 M_2 | O_i | B \rangle &= \sum_j F_j^{B \rightarrow M_1} \int_0^1 dx T_{ij}^I(x) \Phi_{M_1}(x) \\ &+ \int_0^1 d\xi \int_0^1 dx \int_0^1 dy T_{ij}^{II}(\xi, x, y) \Phi_B(\xi) \Phi_{M_1}(x) \Phi_{M_2}(y),\end{aligned}\quad (16)$$

where T_{ij}^I and T_{ij}^{II} denote short-distance interactions and can be calculated perturbatively. $\Phi_X(x)$ are the universal nonperturbative light-cone distribution amplitudes. Using the weak effective Hamiltonian given by Eq.(15), we then obtain the decay amplitudes as:

$$\begin{aligned}A(B^- \rightarrow K_0^{*-} \phi) &= i \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left(a_3 + a_4^p + a_5 - r_\chi^\phi (a_6^p - \frac{1}{2} a_8^p) - \frac{1}{2} (a_7 + a_9 + a_{10}^p) \right)_{K_0^* \phi} \right. \\ &\times \left. 2f_\phi F_1^{BK_0^*}(m_\phi^2) m_{BPc} - f_B f_\phi f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,EW})_{K_0^* \phi} \right\},\end{aligned}\quad (17)$$

$$\begin{aligned}A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \phi) &= i \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left(a_3 + a_4^p + a_5 - r_\chi^\phi (a_6^p - \frac{1}{2} a_8^p) - \frac{1}{2} (a_7 + a_9 + a_{10}^p) \right)_{K_0^* \phi} \right. \\ &\times \left. 2f_\phi F_1^{BK_0^*}(m_\phi^2) m_{BPc} - f_B f_\phi f_{K_0^*} (b_3 - \frac{1}{2} b_{3,EW})_{K_0^* \phi} \right\},\end{aligned}\quad (18)$$

$$\begin{aligned}A(B^- \rightarrow \bar{K}_0^{*0} \rho^-) &= i \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_4^p + r_\chi^{K_0^*} (a_6^p - \frac{1}{2} a_8^p) - \frac{1}{2} a_{10}^p \right)_{\rho K_0^*} \right. \\ &\times \left. 2f_{K_0^*} A_0^{B\rho}(m_{K_0^*}^2) m_{BPc} - f_B f_\rho f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,EW})_{\rho K_0^*} \right\},\end{aligned}\quad (19)$$

$$\begin{aligned}A(B^- \rightarrow K_0^{*-} \rho^0) &= i \frac{G_F}{2} \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} (a_6^p + a_8^p) + a_{10}^p \right)_{\rho K_0^*} \right. \\ &\times \left. 2f_{K_0^*} A_0^{B\rho}(m_{K_0^*}^2) m_{BPc} + \left[a_2 \delta_u^p + \frac{3}{2} (a_9 + a_7) \right]_{K_0^* \rho} 2f_\rho F_1^{BK_0^*}(m_\rho^2) m_{BPc} \right. \\ &\times \left. - f_B f_\rho f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,EW})_{\rho K_0^*} \right\},\end{aligned}\quad (20)$$

$$\begin{aligned}A(\bar{B}^0 \rightarrow K_0^{*-} \rho^+) &= i \frac{G_F}{\sqrt{2}} \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} a_6^p + a_{10}^p + r_\chi^{K_0^*} a_8^p \right)_{\rho K_0^*} \right. \\ &\times \left. 2f_{K_0^*} A_0^{B\rho}(m_{K_0^*}^2) m_{BPc} - f_B f_\rho f_{K_0^*} (b_3 - \frac{1}{2} b_{3,EW})_{\rho K_0^*} \right\},\end{aligned}\quad (21)$$

$$\begin{aligned}
A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0) = & i \frac{G_F}{2} \sum_{p=u,c} \lambda_p^{(s)} \left\{ - \left(-a_4^p - r_\chi^{K_0^*} (a_6^p - \frac{1}{2} a_8^p) + \frac{1}{2} a_{10}^p \right)_{\rho K_0^*} \right. \\
& \times 2f_{K_0^*} A_0^{B\rho}(m_{K_0^*}^2) m_B p_c + \left[a_2 \delta_u^p + \frac{3}{2} (a_9 + a_7) \right]_{K_0^* \rho} 2f_\rho F_1^{BK_0^*}(m_\rho^2) m_B p_c \\
& \left. - f_B f_\rho f_{K_0^*} \left(-b_3 + \frac{1}{2} b_{3,\text{EW}} \right)_{\rho K_0^*} \right\}, \tag{22}
\end{aligned}$$

$$\begin{aligned}
A(B^- \rightarrow K_0^{*-} \omega) = & i \frac{G_F}{2} \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left[a_2 \delta_u^p + 2(a_3 + a_5) + \frac{1}{2} (a_9 + a_7) \right]_{K_0^* \omega} 2f_\omega F_1^{BK_0^*}(m_\omega^2) m_B p_c \right. \\
& - \left(a_1 \delta_u^p + a_4^p + r_\chi^{K_0^*} (a_6^p + a_8^p) + a_{10}^p \right)_{\omega K_0^*} 2f_{K_0^*} A_0^{B\omega}(m_{K_0^*}^2) m_B p_c \\
& \left. - f_B f_\omega f_{K_0^*} (b_2 \delta_u^p + b_3 + b_{3,\text{EW}})_{\omega K_0^*} \right\}, \tag{23}
\end{aligned}$$

$$\begin{aligned}
A(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega) = & i \frac{G_F}{2} \sum_{p=u,c} \lambda_p^{(s)} \left\{ \left[a_2 \delta_u^p + 2(a_3 + a_5) + \frac{1}{2} (a_9 + a_7) \right]_{K_0^* \omega} 2f_\omega F_1^{BK_0^*}(m_\omega^2) m_B p_c \right. \\
& - \left(a_4^p + r_\chi^{K_0^*} (a_6^p - \frac{1}{2} a_8^p) - \frac{1}{2} a_{10}^p \right)_{\omega K_0^*} 2f_{K_0^*} A_0^{B\rho}(m_{K_0^*}^2) m_B p_c \\
& \left. - f_B f_\omega f_{K_0^*} (b_3 - \frac{1}{2} b_{3,\text{EW}})_{\omega K_0^*} \right\}; \tag{24}
\end{aligned}$$

where the ratios r_χ^V and r_χ^S are defined as

$$r_\chi^V(\mu) = \frac{2m_V}{m_b(\mu)} \frac{f_V^\perp(\mu)}{f_V}, \quad r_\chi^S(\mu) = \frac{2m_S^2}{m_b(\mu)(m_2(\mu) - m_1(\mu))}. \tag{25}$$

The order of the arguments of the $a_i^p(M_1 M_2)$ and $b_i(M_1 M_2)$ coefficients is dictated by the subscript $M_1 M_2$, where M_1 shares the same spectator quark with the B meson and M_2 is the emitted meson. For the annihilation part, M_1 is referred to the one containing an anti-quark from the weak vertex, and M_2 contains a quark from the weak vertex. Combining the short-distance nonfactorizable corrections, the effective Wilson coefficients a_i^p have the expressions

$$a_i^p(M_1 M_2) = \left(C_i + \frac{C_{i\pm 1}}{N_c} \right) N_i(M_2) + \frac{C_{i\pm 1}}{N_c} \frac{C_F \alpha_s}{4\pi} \left[V_i(M_2) + \frac{4\pi^2}{N_c} H_i(M_1 M_2) \right] + P_i^p(M_2), \tag{26}$$

where $V_i(M_2)$ account for vertex corrections, $H_i(M_1 M_2)$ for hard spectator interactions and $P_i(M_2)$ for penguin contractions. The coefficients b_i and $b_{i,\text{EW}}$ stand for the contribution of annihilation diagrams.

In QCDF approach, the end-point singularities appear in calculating the twist-3 spectator and annihilation amplitudes. Since the treatment of endpoint divergences is model dependent, sub-leading power corrections generally can be studied only in a phenomenological way. As the most popular way, the end-point divergent integrals are treated as signs of infrared sensitive contributions and parameterized by [32]:

$$\int_0^1 \frac{dy}{y} \rightarrow X_A = (1 + \rho_A e^{i\phi_A}) \ln \frac{m_B}{\Lambda_h}, \quad (27)$$

with the unknown real parameters ρ_A and ϕ_A . More discussion about them will be in Section.IV.

III. THE FAMILY NONUNIVERSAL Z' MODEL

In this section, we will review the main part of the family nonuniversal Z' model briefly. In the current work, for simplicity, we only focus on the models in which the interactions between the Z' boson and fermions are flavor nonuniversal for left-handed couplings and flavor diagonal for right-handed cases. Of course, the analysis can be straightly extended to general cases in which the right-handed couplings are also nonuniversal across generations. The basic formulas of the Z' model with family nonuniversal and/or nondiagonal couplings have been presented in Refs.[19, 23], to which we refer readers for detail. Here, we just review the ingredients needed in this work.

In the gauge basis, the neutral current Lagrangian induced by the Z' boson can be written as

$$\mathcal{L}^{Z'} = -g_2 J'_\mu Z'^\mu, \quad (28)$$

where g_2 is the gauge coupling associated with the additional $U(1)'$ group at the M_W scale. Neglecting the renormalization group (RG) running effect between M_W and $M_{Z'}$ and the mixing between Z' and Z boson of SM, we present the chiral current as

$$J'_\mu = \sum_{i,j} \bar{\psi}_i^I \gamma_\mu [(\varepsilon_{\psi_L})_{ij} P_L + (\varepsilon_{\psi_R})_{ij} P_R] \psi_j^I, \quad (29)$$

where the sum extends over the flavors of fermions, the chirality projection operators are $P_{L,R} \equiv (1 \mp \gamma_5)/2$, the superscript I stands for the weak interaction eigenstates, and ε_{ψ_L} (ε_{ψ_R}) denote the left-handed (right-handed) chiral couplings. ε_{ψ_L} and ε_{ψ_R} are required to be hermitian so as to arrive a real Lagrangian. Accordingly, the mass eigenstates of the chiral fields can be defined

by $\psi_{L,R} = V_{\psi_{L,R}} \psi_{L,R}^I$, and the usual CKM matrix is given by $V_{\text{CKM}} = V_{uL} V_{dL}^\dagger$. Then, the chiral Z' coupling matrices in the physical basis of up-type and down-type quarks are, respectively,

$$B_u^X \equiv V_{uX} \varepsilon_{uX} V_{uX}^\dagger, \quad B_d^X \equiv V_{dX} \varepsilon_{dX} V_{dX}^\dagger \quad (X = L, R). \quad (30)$$

If the ε matrices are not proportional to the identity, the B matrices will have non-zero off-diagonal elements, which induce FCNC interactions at the tree level directly. In this work, we assume that the right-handed couplings are diagonal for simplicity. Thereby, the effective Hamiltonian of the $\bar{b} \rightarrow \bar{s} q \bar{q} (q = u, d)$ transitions mediated by the Z' is

$$\mathcal{H}_{\text{eff}}^{Z'} = \frac{2G_F}{\sqrt{2}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} (\bar{b}s)_{V-A} \sum_q (B_{qq}^L (\bar{q}q)_{V-A} + B_{qq}^R (\bar{q}q)_{V+A}) + \text{h.c.}, \quad (31)$$

where $g_1 = e/(\sin \theta_W \cos \theta_W)$ and $M_{Z'}$ the mass of the new gauge boson. We note the above operators of the forms $(\bar{b}s)_{V-A} (\bar{q}q)_{V-A}$ and $(\bar{b}s)_{V-A} (\bar{q}q)_{V+A}$ already exist in SM, so that we represent the Z' effect as a modification to the Wilson coefficients of the corresponding operators. Hence, we rewrite the eq.(31) as

$$\mathcal{H}_{\text{eff}}^{Z'} = -\frac{G_F}{\sqrt{2}} V_{tb}^* V_{ts} \sum_q \left(\Delta C_3 O_3^{(q)} + \Delta C_5 O_5^{(q)} + \Delta C_7 O_7^{(q)} + \Delta C_9 O_9^{(q)} \right) + \text{h.c.}, \quad (32)$$

where the additional contributions to the SM Wilson coefficients at the M_W scale in terms of Z' parameters are given by

$$\Delta C_{3(5)} = -\frac{2}{3V_{tb}^* V_{ts}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} \left(B_{uu}^{L(R)} + 2B_{dd}^{L(R)} \right) \quad (33)$$

$$\Delta C_{9(7)} = -\frac{4}{3V_{tb}^* V_{ts}} \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 B_{sb}^{L*} \left(B_{uu}^{L(R)} - B_{dd}^{L(R)} \right). \quad (34)$$

Thus we can have a Z' contribution to the QCD penguins $\Delta C_{3(5)}$ as well as the EW penguins $\Delta C_{9(7)}$, in the light of the results found by Buras et al. [42]. In order to show that the new physics is primarily manifest in the EW penguins, we assume $B_{uu}^{L(R)} \simeq -2B_{dd}^{L(R)}$, which have been used widely [25, 27, 28, 30]. As a result, the Z' contributions to the Wilson coefficients at the weak scale are

$$\Delta C_{3(5)} = 0, \quad (35)$$

$$\Delta C_{9(7)} = 4 \frac{|V_{tb}^* V_{ts}|}{V_{tb}^* V_{ts}} \xi^{LL(R)} e^{-i\phi_L}, \quad (36)$$

where

$$\xi^{LX} \equiv \left(\frac{g_2 M_Z}{g_1 M_{Z'}} \right)^2 \left| \frac{B_{sb}^{L*} B_{dd}^X}{V_{tb}^* V_{ts}} \right| \quad (X = L, R), \quad (37)$$

$$\phi_L \equiv \text{Arg}[B_{sb}^L]. \quad (38)$$

Because of the hermiticity of the effective Hamiltonian, the diagonal elements of the effective coupling matrix must be real. However, the off-diagonal elements, such as B_{sb}^L , generally may contain new weak phases. Moreover, the relation $B_{ss}^{L(R)} \simeq B_{dd}^{L(R)}$ follows from the assumptions of universality for the first two families, as required by K and μ decay constraints [23]. Since the major objective of our work is searching for new physics signal, rather than producing acute numerical results, we also assume $B_{qq}^L \simeq B_{qq}^R$, because we expect that $|B_{qq}^L|$ and $|B_{qq}^R|$ should have the same order of magnitude.

It should be emphasized that the other SM Wilson coefficients may also receive contributions from the Z' boson through renormalization group (RG) evolution. With our assumption that no significant RG running effect between $M_{Z'}$ and M_W scales, the RG evolution of the modified Wilson coefficients is exactly the same as the ones in SM [42]. The numerical results of Wilson coefficients in the naive dimensional regularization (NDR) scheme at the scale $\mu = 2.1\text{GeV}$ ($\mu_h = 1\text{GeV}$) are listed in Table I for convenience.

In summary, we list here our simplifications to a general Z' model: we assume (i) no right-handed flavor-changing couplings ($B_{ij}^R = 0$ for $i \neq j$), (ii) no significant RG running effect between $M_{Z'}$ and M_W scales, (iii) negligible Z' effect on the QCD penguin ($\Delta C_{3,5} = 0$) so that the new physics is manifestly isospin violating, (iv) $|B_{qq}^L|$ and $|B_{qq}^R|$ are same so as to reduce the number of parameters. With these simplifications, we have only two parameters left in the model. So, this approach provides a minimal way to introduce the Z' effect in the concerned decay modes. Of course, more general Z' models are possible.

Now, the only task left is to constraint the parameters within the existing experimental data. Generally, $g_2/g_1 \sim 1$ is expected, if both the $U(1)$ gauge groups have the same origin from some grand unified theories. We also hope $M_Z/M_{Z'} \sim 0.1$ so that TeV scale neutral Z' boson could be detected at LHC. Theoretically, one can fit the left three parameters $|B_{sb}^L|$, $|B_{dd}^X|$ and new weak phase ϕ_L with the accurate data from B factories and other experiments such as Tavatron and LHC. For example, B_{sb}^L and ϕ_L could be extracted from B_s - \bar{B}_s mixing as well as $B \rightarrow K^{(*)} \ell^+ \ell^-$ decays. To resolve the mass difference between B_s and \bar{B}_s , $|B_{sb}^L| \sim |V_{tb} V_{ts}^*|$ is required [25, 28, 43]. In Refs.[28], the authors got the ϕ_L is about -80° by fitting data of $B_s - \bar{B}_s$ mixing and $B \rightarrow K^{(*)} l^+ l^-$

TABLE I: The Wilson coefficients C_i within SM and with the contribution from Z' boson included in NDR scheme at the scale $\mu = 2.1$ GeV and $\mu_h = 1.0$ GeV.

Wilson coefficients	$\mu = 2.1$ GeV		$\mu_h = 1.0$ GeV	
	C_i^{SM}	$\Delta C_i^{Z'}$	C_i^{SM}	$\Delta C_i^{Z'}$
C_1	1.135	0	1.224	0
C_2	-0.283	0	-0.429	0
C_3	0.021	$0.09\xi^{LL} - 0.02\xi^{LR}$	0.034	$0.15\xi^{LL} - 0.04\xi^{LR}$
C_4	-0.049	$-0.20\xi^{LL} + 0.01\xi^{LR}$	-0.072	$-0.31\xi^{LL} + 0.03\xi^{LR}$
C_5	0.010	$0.03\xi^{LL} + 0.02\xi^{LR}$	0.010	$0.02\xi^{LL} + 0.02\xi^{LR}$
C_6	-0.06	$-0.26\xi^{LL} + 0.03\xi^{LR}$	-0.104	$-0.44\xi^{LL} + 0.07\xi^{LR}$
C_7/α_{em}	-0.018	$5.3\xi^{LL} - 461\xi^{LR}$	-0.023	$6.3\xi^{LL} - 457\xi^{LR}$
C_8/α_{em}	0.081	$2.43\xi^{LL} - 286\xi^{LR}$	0.134	$4.8\xi^{LL} - 497\xi^{LR}$
C_9/α_{em}	-1.266	$-594\xi^{LL} + 6.1\xi^{LR}$	-1.366	$-643\xi^{LL} + 7.8\xi^{LR}$
C_{10}/α_{em}	0.321	$178\xi^{LL} - 1.0\xi^{LR}$	0.483	$257\xi^{LL} - 1.9\xi^{LR}$
$C_{7\gamma}$	-0.345	—	-0.395	—
C_{8g}	-0.161	—	-0.181	—

decays. Subsequently, with B_{sb}^L and ϕ_L arrived and experimental data of $B \rightarrow \pi\pi, K\pi, K\rho$ and $K^{(*)}\phi$, B_{qq}^L and B_{qq}^R could be extracted analogously. Specifically, the CP asymmetries in $B \rightarrow K\phi, K\pi$ can be resolved if $|B_{sb}^L B_{ss}^{L,R}| \sim |V_{tb} V_{ts}^*|$, which indicates $|B_{qq}^{L,R}| \sim 1$. However, we have one remark here. In dealing with the nonleptonic B decays, because different groups used different factorization approach, the fitted results are different, but all results have same order. Noted that the detailed constraint of these parameters is beyond the scope of current work and can be found in many references [27, 28]. Summing up above analysis, we thereby assume that $\xi = \xi^{LL} = \xi^{LR} \in (10^{-3}, 10^{-2})$ and $\phi_L \in (-60^\circ, -90^\circ)$ so as to prob the new physics effect for maximum range.

IV. NUMERICAL RESULTS AND DISCUSSION

In this section, to begin with, we will give the parameters used in this work. Since it is not clear whether the scalar meson $K_0^*(1430)$ belongs to the first orbital excited state (S1) or the low lying

resonance (S2), we will calculate the processes under both scenarios. In the calculation, the decay constants and Gegenbauer moments obtained within the QCD sum rules method under different scenarios are presented as follows [3]:

$$\begin{aligned} \mathbf{S1}: \bar{f}_{K_0^*}(1.0\text{GeV}) &= -300\text{MeV}; \bar{f}_{K_0^*}(2.1\text{GeV}) = -370\text{MeV}; B_1(1.0\text{GeV}) = 0.58; \\ B_1(2.1\text{GeV}) &= 0.39; B_3(1.0\text{GeV}) = -1.20; B_3(2.1\text{GeV}) = -0.70; \end{aligned} \quad (39)$$

$$\begin{aligned} \mathbf{S2}: \bar{f}_{K_0^*}(1.0\text{GeV}) &= 445\text{MeV}; \bar{f}_{K_0^*}(2.1\text{GeV}) = 550\text{MeV}; B_1(1.0\text{GeV}) = -0.57; \\ B_1(2.1\text{GeV}) &= -0.39; B_3(1.0\text{GeV}) = -0.42; B_3(2.1\text{GeV}) = -0.25. \end{aligned} \quad (40)$$

In QCD sum rules method, the major parameter is the Borel window, which takes large uncertainty to the parameters listed above. In Ref.[3], the authors had discussed the errors caused by them in great detail and found that $B_{1,3}$ will take 30% changes. As a result, we will not discuss this part any more in the current work.

For the vector mesons, the longitudinal and transverse decay constants are list as:

$$\begin{aligned} f_\rho &= 216 \text{ MeV}, & f_\omega &= 187 \text{ MeV}, & f_\phi &= 215 \text{ MeV}, \\ f_\rho^\perp &= 165 \text{ MeV}, & f_\omega^\perp &= 151 \text{ MeV}, & f_\phi^\perp &= 186 \text{ MeV}, \end{aligned} \quad (41)$$

where the values are taken from [44]. In the LCADs of vectors, the Gegenbauer moments α_n^V and $\alpha_{n,\perp}^V$ have been studied within the QCD sum rule method. Here, we will employ the most recent updated values [45]

$$\alpha_2^{\rho,\omega} = 0.15, \quad \alpha_{2,\perp}^{\rho,\omega} = 0.14, \quad \alpha_2^\phi = 0.18, \quad \alpha_{2,\perp}^\phi = 0.14, \quad (42)$$

and $\alpha_1^V = 0, \alpha_{1,\perp}^V = 0$.

As stated earlier, various form factors for $B \rightarrow S, V$ transitions have been evaluated in cLFQM [38]. In this model form factors are first calculated in the spacelike region and their momentum dependence is fitted to a 3-parameter form

$$F(q^2) = \frac{F(0)}{1 - a(q^2/m_B^2) + b(q^2/m_B^2)^2}. \quad (43)$$

The parameters a, b and $F(0)$ relevant for our purposes are summarized in Table.II.

In Refs.[3, 4, 14], it was found that in decay modes with scalars the main theoretical uncertainties are due to the weak annihilations, especially for the penguin dominated ones. In $B \rightarrow PP, PV$ decays, the annihilation amplitudes are helicity suppressed because the helicity of one of final

states cannot match with that of its quarks. However, this helicity suppression can be alleviated in the decay modes with scalar because of nonvanishing orbital angular momentum. Thus, weak annihilation contribution to $B \rightarrow SP(V)$ is much larger than the $B \rightarrow PP(V)$ case. However, as stated before, the end-point singularity appears in calculating the annihilation contribution, and then two free parameters, ρ_A and ϕ_A , are introduced phenomenally. In Ref. [4], it is found that the behavior of SV is similar to the longitudinal part of VV . Fortunately, with experimental data, it presents the moderate value of nonuniversal annihilation phase $\phi_A = -40^\circ$ for $B \rightarrow VV$ decay modes [32]. Therefore, for $B \rightarrow SV$, we conservatively take $\phi_A = (-40 \pm 20)^\circ$ with $\rho_A = 0.6 \pm 0.2$, which also assures that the hadronic uncertainties are considerably reduced. Furthermore, the endpoint divergence X_H in the hard spectator contributions can also be parameterized in the same manner.

Within above parameters and formulas, we calculate the branching fractions of these decays in SM and the family nonuniversal Z' model under two different scenarios. Together with partial experimental results, the results under are exhibited in Table.III, respectively. For the center values, we adopt $\xi = 0.005$ and $\phi_L^{sb} = -80^\circ$. For all theoretical predictions, the first errors arise from the power corrections of weak annihilation and hard spectator interactions characterized by the parameters $X_{A,H}$. To obtain the second errors of the Z' model results, we scan randomly the points in their own possible parameter spaces.

Comparing our predictions of SM with those in Ref. [6] (considering the typos), there are few differences. Some reasons are list as follows: (1) In the Ref. [6], for the parameterizations of singularities, the center values correspond to $\rho_{A,H} = 0$ and $\phi_{A,H} = 0$, while we set $\rho_{A,H} = 0.6$ and $\phi_{A,H} = -40^\circ$; (2) The difference of Wilson coefficients, caused by the top quark mass and other part parameters, will change the results slightly; (3) In this work, the different form factors of $B \rightarrow K_0^*(1430)$ are used under different scenarios, but they adopted same values in the ref.[4].

TABLE II: Form factors of $B \rightarrow \rho, K_0^*(1430)$ transitions obtained in the covariant light-front model [38].

F	$F(0)$	$F(q_{\max}^2)$	a	b	F	$F(0)$	$F(q_{\max}^2)$	a	b
$V^{B\rho}$	0.27	0.79	1.84	1.28	$A_0^{B\rho}$	0.28	0.76	1.73	1.20
$A_1^{B\rho}$	0.22	0.53	0.95	0.21	$A_2^{B\rho}$	0.20	0.57	1.65	1.05
$F_1^{BK_0^*}[\text{S1}]$	0.21	0.52	1.59	0.91	$F_0^{BK_0^*}[\text{S1}]$	0.21	0.30	0.59	0.09
$F_1^{BK_0^*}[\text{S2}]$	0.26	0.70	1.52	0.64	$F_0^{BK_0^*}[\text{S2}]$	0.26	0.33	0.44	0.05

TABLE III: Branching fractions (in units of 10^{-6}) under the different Scenarios

Decay Mode	S1		S2		Expt
	SM	SM+Z'	SM	SM+Z'	
$B^- \rightarrow K_0^{*-} \phi$	$2.4^{+5.1}_{-1.8}$	$3.8^{+5.6+1.8}_{-2.5-1.2}$	$22.6^{+19.7}_{-8.6}$	$16.5^{+18.5+5.0}_{-10.3-7.9}$	$7.0 \pm 1.3 \pm 0.9$
$\bar{B}^0 \rightarrow K_0^{*0} \phi$	$2.2^{+4.9}_{-1.7}$	$4.7^{+5.2+4.4}_{-2.6-2.2}$	$22.4^{+19.4}_{-8.4}$	$21.2^{+18.9+3.0}_{-8.4-4.6}$	$4.3 \pm 0.6 \pm 0.4$
$B^- \rightarrow \bar{K}_0^{*0} \rho^-$	$11.7^{+8.4}_{-4.0}$	$11.5^{+7.7+3.6}_{-4.7-3.8}$	$45.5^{+20.6}_{-10.4}$	$41.4^{+19.7+4.4}_{-11.9-5.7}$	27 ± 5.5
$B^- \rightarrow \bar{K}_0^{*-} \rho^0$	$7.2^{+4.5}_{-2.3}$	$18.2^{+5.9+27.4}_{-3.4-10.0}$	$17.6^{+9.0}_{-4.4}$	$15.9^{+8.4+6.9}_{-6.1-5.6}$	
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0$	$4.6^{+1.7}_{-0.7}$	$3.9^{+1.4+6.4}_{-0.8-1.8}$	$24.5^{+7.0}_{-3.8}$	$33.3^{+7.7+32.1}_{-4.8-8.5}$	
$\bar{B}^0 \rightarrow \bar{K}_0^{*-} \rho^+$	$10.7^{+8.5}_{-3.7}$	$14.4^{+9.2+8.7}_{-5.1-4.1}$	$44.7^{+21.3}_{-10.5}$	$54.1^{+22.4+25.7}_{-12.6-8.6}$	
$B^- \rightarrow \bar{K}_0^{*-} \omega$	$3.6^{+3.1}_{-1.4}$	$7.8^{+3.7+10.1}_{-1.7-3.8}$	$12.6^{+7.2}_{-3.5}$	$13.7^{+7.5+4.8}_{-3.9-1.4}$	28 ± 12
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega$	$3.9^{+1.3}_{-0.6}$	$4.0^{+1.3+1.2}_{-0.7-0.6}$	$10.6^{+4.4}_{-2.1}$	$10.7^{+4.4+5.4}_{-2.7-2.6}$	

In Tables.III, for $K_0^*(1430)\phi$ channels, though the central values of the predicted under S1 are smaller than the experimental data, they are accommodated with the large uncertainties. However, in S2, the theoretical results are much larger than the data, and cannot agree with data even with uncertainties. It should be noted that in this work we have not included the errors from uncertainties of meson distribution amplitudes (B_1 and B_3). Even with those uncertainties, the theoretical results are still larger than the upper limits of the data. These theoretical results also agree with the results from pQCD approach [8]. For $\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0$ channels, contrary to $K_0^*(1430)\phi$, the result of S2 agree with data well and the prediction of S1 is much smaller than the data. Since for $\bar{B}^0 \rightarrow \bar{K}_0^{*-} \rho^+$ there is large uncertainty in the experimental data, the theoretical results under both S1 and S2 can accommodate the data with large uncertainties theoretically. That's to say, it is impossible to explain all data under one settled scenario simultaneously.

When adding the contribution of the Z' gauge boson, as shown in the table, the Z' gauge boson changes the branching fractions under both two different scenarios. For $B \rightarrow K_0^* \phi$ channels dominated by the weak annihilation, the Z' will enhance the branching fractions in S1, while in S2 the branching ratios are decreased. The reason is that the weak annihilation is proportional to the decay constant $f_{K_0^*}$, which has different sign in different scenarios. For $B^- \rightarrow \bar{K}_0^{*0} \rho^-$ and $\bar{B}^0 \rightarrow \bar{K}_0^{*-} \rho^+$, as the scalar particle is the emitted particle, the whole amplitudes are proportional to the decay constant $f_{K_0^*}$, thus the new physics contribution have same behavior in different sce-

narios. For channels with ρ^0 or ω , the spectator quarks enters not only the scalars but also the vectors, the amplitudes become more complicate, and we cannot describe the relation between new physics and branching fractions apparently.

Compared to the experimental data, the Z' boson could change the branching fractions remarkably and alleviate the disparities. However, we cannot achieve a definite conclusion yet whether K_0^* belongs to the ground states or the first orbital excited states. Moreover, for most modes except $B^- \rightarrow \bar{K}_0^{*0-} \rho^0(\omega)$, the new physics contribution might be clouded by the uncertainties taken by the weak annihilations. Thus, it is also very difficult to search for Z' effect in these decays. Specifically, for decays $B^- \rightarrow \bar{K}_0^{*-} \rho^0(\omega)$, Z' boson could enhance the branching fractions more than 2 times, we hope these two channels could be measured in the LHC or Super-b factories in future so as to probe the Z' gauge boson.

To test the isospin symmetry and prob new physics, we define two ratios:

$$R_1 = \frac{Br(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0)}{Br(\bar{B}^0 \rightarrow \bar{K}_0^{*-} \rho^+)} = 0.96_{-0.43}^{+1.07} \quad [\text{Exp.}]; \quad (44)$$

$$R_2 = \frac{\tau(\bar{B}^0)}{\tau(B^-)} \cdot \frac{Br(B^- \rightarrow \bar{K}_0^{*-} \phi)}{Br(\bar{B}^0 \rightarrow \bar{K}_0^{*0} \phi)} = 1.52_{-0.51}^{+0.71} \quad [\text{Exp.}], \quad (45)$$

where the experimental results are also given and all uncertainties are added in quadrature. In the isospin limit, $R_1 = 1/2$ and $R_2 = 1$ are expected to hold. Here, we list the theoretical results under different scenarios in different models:

$$R_1[\text{SM}] = \begin{cases} 0.43_{-0.10}^{+0.13}, & \text{S1;} \\ 0.55_{-0.08}^{+0.08}, & \text{S2.} \end{cases}; \quad R_1[\text{SM} + Z'] = \begin{cases} 0.27_{-0.08-0.17}^{+0.15+0.31}, & \text{S1;} \\ 0.62_{-0.09-0.07}^{+0.09+0.20}, & \text{S2.} \end{cases} \quad (46)$$

$$R_2[\text{SM}] = \begin{cases} 1.00_{-0.04}^{+0.04}, & \text{S1;} \\ 0.94_{-0.01}^{+0.00}, & \text{S2.} \end{cases}; \quad R_2[\text{SM} + Z'] = \begin{cases} 0.76_{-0.18-0.18}^{+0.16+0.24}, & \text{S1;} \\ 0.73_{-0.27-0.24}^{+0.16+0.19}, & \text{S2.} \end{cases}; \quad (47)$$

In the above results, the theoretical uncertainties are reduced since they are ratios of branch fractions. We see that the symmetries are almost held in SM. However, the data shows that the isospin symmetries are violated, which means that the large weak annihilation may break the isospin symmetry remarkably. When adding Z' contribution, except R_1 under S2, the isospin symmetries are broken in an opposite direction. However, the family nonuniversal Z' model cannot be ruled out due to large uncertainties in the experiments.

Finally, we will discuss the CP asymmetries of these decays. For the charged mode $B^- \rightarrow K_0^{*-} \phi$, because $|V_{ub}V_{us}|(\lambda^4) \ll |V_{tb}V_{ts}|(\lambda^2)$ and there is no tree contribution in the neutral mode

TABLE IV: The direct CP asymmetry (%) under the different scenarios

Decay Mode	S1		S2	
	SM	SM+ Z'	SM	SM+ Z'
$B^- \rightarrow \bar{K}_0^{*0} \rho^-$	6_{-2}^{+4}	6_{-3-1}^{+6+1}	2_{-1}^{+2}	2_{-1-0}^{+2+0}
$B^- \rightarrow \bar{K}_0^{*-} \rho^0$	4_{-3}^{+4}	-3_{-1-2}^{+2+4}	-1_{-4}^{+3}	6_{-3-6}^{+3+10}
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho^0$	9_{-38}^{+26}	24_{-42-12}^{+52+21}	-11_{-13}^{+10}	-9_{-11-2}^{+9+4}
$\bar{B}^0 \rightarrow \bar{K}_0^{*-} \rho^+$	1_{-2}^{+1}	-2_{-1-0}^{+1+1}	1_{-0}^{+0}	1_{-0-0}^{+0+0}
$B^- \rightarrow \bar{K}_0^{*-} \omega$	3_{-7}^{+6}	-4_{-3-3}^{+2+5}	-1_{-5}^{+4}	-4_{-4-6}^{+3+5}
$\bar{B}^0 \rightarrow \bar{K}_0^{*0} \omega$	16_{-39}^{+26}	17_{-40-5}^{+28+7}	-19_{-15}^{+14}	-4_{-19-13}^{+15+19}

$\bar{B}^0 \rightarrow K_0^{*0} \phi$, the direct CP asymmetries are almost zero in both SM and the Z' model. For $B \rightarrow K_0^* \rho$, although the CKM elements are suppressed, the tree operators with large Wilson coefficients appear in the emission diagrams, so the amplitudes of tree and penguin may have comparable magnitudes. Thus, large CP asymmetries in these decays are expected, just like decays $B \rightarrow K\pi$ and $B \rightarrow K\rho$. In Table.IV, we give the CP asymmetries of $B \rightarrow K_0^* \rho$ in both SM and the concerned new physics model under different scenarios. From the table, we firstly note that $\bar{B}^0 \rightarrow \bar{K}_0^{*0} \rho(\omega)$ have large asymmetries, and different scenarios have different signs but with large uncertainties. If we can calculate the annihilation accurately within some effective approach in future, this parameter could be used to distinguish the scenarios. Secondly, for $B^- \rightarrow \bar{K}_0^{*-} \rho^0(\omega)$, the Z' could change the signs of the center values, and these two decays can be used in probing new physics effect.

V. SUMMARY

Motivated by recent measurements of decays $B \rightarrow K_0^* \rho$ and $K_0^* \phi$, we studied the branching fractions of these decays both in SM and in the family nonuniversal Z' model within the QCDF framework. Because it is not clear whether K_0^* is the lying state or the first orbital excited state, we calculate them under two different scenarios. For these decay modes with scalar meson, the weak annihilations play more important roles than that in $B \rightarrow PP$ and PV decays, so that they will take large uncertainties. From this point, an effective way that could calculate the annihilations reliably is needed. Comparing with the experimental results, we found different channels favor different scenarios. Moreover, in order to account for the large isospin asymmetries in the data,

large weak annihilations are also required. Adding the contribution of the family nonuniversal Z' boson, we note that both the branching fractions and their ratios are changed remarkably. However, we cannot identify the character of the scalar meson K_0^* , either. Furthermore, for most channels, the Z' contribution will be buried by large uncertainties, except for decays $B^- \rightarrow \bar{K}_0^{*-} \rho^0(\omega)$.

In this work, we also calculated the CP asymmetries of these decays and found the CP asymmetries of $B \rightarrow K_0^* \phi$ are almost zero. In different scenarios, the CP asymmetries of $B^- \rightarrow \bar{K}_0^{*-} \rho^0(\omega)$ have different signs, thus they can be used to classify the scalar K_0^* . If its character is identified, we accordingly could use these results to probe the new gauge boson Z' , because it changes the signs of CP asymmetries. All above results could be tested in the running LHCb or the Super-b factories in future.

Acknowledgement

The work of Y.Li is supported by the National Science Foundation (Nos.11175151) and the Natural Science Foundation of Shandong Province (ZR2010AM036).

-
- [1] H. -Y. Cheng and J. G. Smith, Ann. Rev. Nucl. Part. Sci. **59**, 215 (2009) [arXiv:0901.4396 [hep-ph]].
 - [2] P. Minkowski and W. Ochs, Eur. Phys. J. C **9**, 283 (1999) [arXiv:hep-ph/9811518].
 - [3] H. Y. Cheng, C. K. Chua and K. C. Yang, Phys. Rev. D **73**, 014017 (2006) [arXiv:hep-ph/0508104].
 - [4] H. Y. Cheng, C. K. Chua and K. C. Yang, Phys. Rev. D **77**, 014034 (2008) [arXiv:0705.3079 [hep-ph]].
 - [5] C. H. Chen, C. Q. Geng, Y. K. Hsiao and Z. T. Wei, Phys. Rev. D **72**, 054011 (2005) [arXiv:hep-ph/0507012].
 - [6] C. -H. Chen and C. -Q. Geng, Phys. Rev. D **75**, 054010 (2007) [hep-ph/0701023].
 - [7] Y. L. Shen, W. Wang, J. Zhu and C. D. Lu, Eur. Phys. J. C **50**, 877 (2007) [arXiv:hep-ph/0610380];
 - [8] C. S. Kim, Y. Li and W. Wang, Phys. Rev. D **81**, 074014 (2010) [arXiv:0912.1718 [hep-ph]];
 - [9] X. Liu, Z. -Q. Zhang and Z. -J. Xiao, Chin. Phys. C **34**, 157 (2010) [arXiv:0904.1955 [hep-ph]].
 - [10] X. Liu and Z. -J. Xiao, Commun. Theor. Phys. **53**, 540 (2010) [arXiv:1004.0749 [hep-ph]].
 - [11] Z. -Q. Zhang, Phys. Rev. D **82**, 034036 (2010) [arXiv:1006.5772 [hep-ph]].

- [12] Z. -Q. Zhang, Phys. Rev. D **82**, 114016 (2010) [arXiv:1106.0103 [hep-ph]].
- [13] X. Liu, Z. -J. Xiao and Z. -T. Zou, arXiv:1105.5761 [hep-ph].
- [14] Y. Li, X. -J. Fan, J. Hua and E. -L. Wang, Phys. Rev. D **85**, 074010 (2012) [arXiv:1111.7153 [hep-ph]].
- [15] J. P. Lees *et al.* [The BABAR Collaboration], arXiv:1112.3896 [hep-ex].
- [16] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **98**, 051801 (2007) [hep-ex/0610073].
- [17] B. Aubert *et al.* [BABAR Collaboration], Phys. Rev. Lett. **101**, 161801 (2008) [arXiv:0806.4419 [hep-ex]].
- [18] Y. Gao, PhD thesis, AAT-3356872.
- [19] For recent review, see P. Langacker, Rev. Mod. Phys. **81**, 1199 (2009) [arXiv:0801.1345 [hep-ph]].
- [20] E. Nardi, Phys. Rev. D **48** (1993) 1240 [hep-ph/9209223];
J. Bernabeu, E. Nardi and D. Tommasini, Nucl. Phys. B **409** (1993) 69 [hep-ph/9306251];
V. Barger, M. S. Berger and R. J. N. Phillips, Phys. Rev. D **52** (1995) 1663 [hep-ph/9503204];
M. B. Popovic and E. H. Simmons, Phys. Rev. D **62** (2000) 035002 [hep-ph/0001302];
T. G. Rizzo Phys. Rev. D **59** (1998) 015020 [hep-ph/9806397].
- [21] S. Chaudhuri, S. -W. Chung, G. Hockney and J. D. Lykken, Nucl. Phys. B **456**, 89 (1995) [hep-ph/9501361].
- [22] R. W. Robinett and J. L. Rosner, Phys. Rev. D **25**, 3036 (1982) [Erratum-ibid. D 27, 679 (1983)].
- [23] P. Langacker and M. Plumacher, Phys. Rev. D **62**, 013006 (2000) [arXiv:hep-ph/0001204] and references therein.
- [24] V. Barger, *et. al.*, Phys. Lett. B **580**, 186 (2004) [arXiv:hep-ph/0310073];
V. Barger, *et. al.*, Phys. Lett. B **598**, 218 (2004) [arXiv:hep-ph/0406126].
- [25] V. Barger, *et. al.*, Phys. Rev. D **80**, 055008 (2009) [arXiv:0902.4507 [hep-ph]];
V. Barger, *et. al.*, JHEP **0912**, 048 (2009) [arXiv:0906.3745 [hep-ph]].
- [26] A. K. Alok, S. Baek and D. London, JHEP **1107**, 111 (2011) [arXiv:1010.1333 [hep-ph]];
H. D. Kim, S. -G. Kim and S. Shin, arXiv:1205.6481 [hep-ph].
- [27] K. Cheung, *et. al.*, Phys. Lett. B **652**, 285 (2007) [arXiv:hep-ph/0604223];
C. W. Chiang, *et. al.*, JHEP **0608**, 075 (2006) [arXiv:hep-ph/0606122];
- [28] Q. Chang, X. Q. Li and Y. D. Yang, JHEP **0905**, 056 (2009) [arXiv:0903.0275 [hep-ph]];
Q. Chang, X. Q. Li and Y. D. Yang, JHEP **1002**, 082 (2010) [arXiv:0907.4408 [hep-ph]]
Q. Chang, X. Q. Li and Y. D. Yang, JHEP **1004**, 052 (2010) [arXiv:1002.2758 [hep-ph]].

- [29] J. Hua, C. S. Kim and Y. Li, Phys. Lett. B **690**, 508 (2010) [arXiv:1002.2532 [hep-ph]];
 J. Hua, C. S. Kim and Y. Li, Eur. Phys. J. C **69**, 139 (2010) [arXiv:1002.2531 [hep-ph]]
 Q. Chang and Y. H. Gao, Nucl. Phys. B **845**, 179 (2011) [arXiv:1101.1272 [hep-ph]];
 Y. Li, J. Hua and K. C. Yang, Eur. Phys. J. C **71**, 1775 (2011) [arXiv:1107.0630 [hep-ph]];
 X. -Q. Li, Y. -M. Li, G. -R. Lu and F. Su, JHEP **1205**, 049 (2012) [arXiv:1204.5250 [hep-ph]].
- [30] A. Arhrib, K. Cheung, C. -W. Chiang and T. -C. Yuan, Phys. Rev. D **73**, 075015 (2006) [hep-ph/0602175].
- [31] C. -W. Chiang, Y. -F. Lin and J. Tandean, JHEP **1111**, 083 (2011) [arXiv:1108.3969 [hep-ph]].
- [32] M. Beneke, G. Buchalla, M. Neubert and C. T. Sachrajda, Phys. Rev. Lett. **83**, 1914 (1999) [arXiv:hep-ph/9905312];
 M. Beneke, G. Buchalla, M. Neubert and C. T. Sachrajda, Nucl. Phys. B **591**, 313 (2000) [arXiv:hep-ph/0006124];
 M. Beneke and M. Neubert, Nucl. Phys. B **675**, 333 (2003) [arXiv:hep-ph/0308039];
 M. Beneke, J. Rohrer and D. Yang, Nucl. Phys. B **774**, 64 (2007) [hep-ph/0612290].
- [33] M. -Z. Yang, Phys. Rev. D **73**, 034027 (2006) [Erratum-ibid. D **73**, 079901 (2006)] [hep-ph/0509103].
- [34] T. M. Aliev, K. Azizi and M. Savci, Phys. Rev. D **76**, 074017 (2007) [arXiv:0710.1508 [hep-ph]].
- [35] Y. -M. Wang, M. J. Aslam and C. -D. Lu, Phys. Rev. D **78**, 014006 (2008) [arXiv:0804.2204 [hep-ph]].
- [36] Y. -J. Sun, Z. -H. Li and T. Huang, Phys. Rev. D **83**, 025024 (2011) [arXiv:1011.3901 [hep-ph]].
- [37] R. -H. Li, C. -D. Lu, W. Wang and X. -X. Wang, Phys. Rev. D **79**, 014013 (2009) [arXiv:0811.2648 [hep-ph]].
- [38] H. -Y. Cheng, C. -K. Chua and C. -W. Hwang, Phys. Rev. D **69**, 074025 (2004) [hep-ph/0310359].
- [39] P. Ball and R. Zwicky, Phys. Rev. D **71**, 014029 (2005) [hep-ph/0412079].
- [40] C. D. Lu, Y. M. Wang and H. Zou, Phys. Rev. D **75**, 056001 (2007) [arXiv:hep-ph/0612210].
- [41] M. Wirbel, S. Stech, and M. Bauer, Z. Phys. C **29**, 637 (1985);
 M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **34**, 103 (1987);
 M. Bauer, B. Stech, and M. Wirbel, Z. Phys. C **42**, 671 (1989).
- [42] G. Buchalla, A. J. Buras, and M. E. Lautenbacher, Rev. Mod. Phys. **68**, 1125 (1996) [hep-ph/9512380].
- [43] A. K. Alok, S. Baek and D. London, JHEP **1107**, 111 (2011) [arXiv:1010.1333 [hep-ph]].

- [44] P. Ball, G. W. Jones and R. Zwicky, Phys. Rev. D **75**, 054004 (2007) [hep-ph/0612081].
- [45] P. Ball and G. W. Jones, JHEP **0703**, 069 (2007) [hep-ph/0702100 [HEP-PH]].